UCRHEP-T149
July 1995

Neutrino Masses in an Extended Gauge Model with E_6 Particle Content

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Abstract

Naturally light singlet neutrinos which mix with the usual doublet neutrinos are possible if the supersymmetric standard gauge model is extended to include a specific additional $U(1)$ factor derivable from an E_6 decomposition. The low-energy particle content of the model is limited to the fundamental **27** representations of E_6.

The three known neutrinos ν_e, ν_μ, and ν_τ are each a component of an $SU(2) \times U(1)$ doublet, pairing with the left-handed projections of the charged leptons e, μ, and τ respectively. They are generally considered to be Majorana fermions with very small masses arising from the so-called "seesaw" mechanism.[1] This means that there should be three heavy neutral fermion singlets $N_{1,2,3}$ which also couple to $\nu_{e,\mu,\tau}$ through the usual Higgs doublet $\Phi = (\phi^+, \phi^0)$ of the standard model. As ϕ^0 acquires a nonzero *vev* (vacuum expectation value), a Dirac mass term m_D linking ν and N is obtained, yielding the well-known result $m_\nu \simeq m_D^2/m_N$. The most natural origin of $N_{1,2,3}$ is that associated with a left-right model where they can be identified as the right-handed counterparts of the left-handed neutrinos. As the $SU(2)_L \times SU(2)_R \times U(1)$ gauge symmetry breaks down to the standard $SU(2) \times U(1)$, a large m_N may be obtained.

Whereas the usual seesaw mechanism is based on a 2×2 matrix

$$\mathcal{M}_2 = \begin{pmatrix} 0 & m_D \\ m_D & m_N \end{pmatrix} \tag{1}$$

with the doublet neutrino getting a small mass, there is also the simple variation where it is the <u>singlet</u> neutrino which gets a small mass. Consider the left-handed fermion doublets (ν_E, E) and (E^c, N_E^c) transforming as $(2, -1/2)$ and $(2,1/2)$ respectively under the standard $SU(2) \times U(1)$. Add a neutral fermion singlet S and forbid it to have a Majorana mass term by an appropriate symmetry. The 3×3 mass matrix spanning ν_E, N_E^c, and S may then be given by

$$\mathcal{M}_3 = \begin{pmatrix} 0 & m_E & m_1 \\ m_E & 0 & m_2 \\ m_1 & m_2 & 0 \end{pmatrix}, \tag{2}$$

where $m_{1,2}$ are proportional to the *vev* of an appropriate Higgs doublet and m_E is now an allowed gauge-invariant mass. For $m_{1,2} << m_E$, we then have $m_S \simeq 2m_1m_2/m_E$. If $\mathcal{M}_3$ is also linked to $\mathcal{M}_2$, then the light singlet S will also mix with the usual doublet neutrinos.

If a light singlet neutrino exists in addition to the three doublet neutrinos, a compre-

hensive picture of neutrino oscillations and hot dark matter becomes possible.[2] This is especially so because of the recent results of the LSND (Liquid Scintillator Neutrino Detector) experiment[3] which may be interpreted as ν_μ oscillating to ν_e with a Δm^2 of a few eV^2. To avoid the severe constraint on the effective number of neutrinos from big-bang nucleosynthesis,[4] the singlet neutrino may be used only to account for the solar data by mixing with ν_e in the matter-enhanced small-angle solution or the long-wavelength large-angle solution.

A good model for a light singlet neutrino should have an appropriate symmetry which forbids it to have a Majorana mass term, as already noted. It is of course easy to impose such a symmetry, but for it to be natural, it should come from a more fundamental framework, such as grand unification or string theory for example. As it turns out, Eq. (2) is a natural consequence of the superstring-inspired E_6 model.[5] Unfortunately, the corresponding m_N of Eq. (1) is zero there.[6] This means that $\nu_{e,\mu,\tau}$ combine with $N_{1,2,3}$ to form Dirac neutrinos and their small masses are unexplained. On the other hand, gravitationally induced nonrenormalizable interactions[7] may produce large Majorana mass terms for both N and S, in which case $\nu_{e,\mu,\tau}$ are again naturally light by virtue of the seesaw mechanism, but they will be the only ones.

The low-energy gauge symmetry of a superstring-inspired E_6 model is often taken to be $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_\eta$, because the flux mechanism of symmetry breaking in string theory involves the adjoint representation and it is not possible[5] to break E_6 all the way down to the gauge symmetry of the standard model. If only one extra $U(1)$ factor is present, then it is necessarily $U(1)_\eta$, according to which both N and S transform nontrivially. They are thus protected by this gauge symmetry from acquiring large Majorana masses. For the nonrenormalizable mechanism of Ref. [7] to work, the $U(1)_\eta$ must also be broken at an intermediate scale by *vev*'s along the N and S directions. To obtain a light neutrino doublet

with Eq. (1) as well as a light neutrino singlet with Eq. (2), the idea then is to replace $U(1)_\eta$ with another $U(1)$ under which N is trivial but S is not, so that only the former may acquire a large Majorana mass. In the following this extended gauge model is described.

Consider the maximal subgroup $SU(3)_C \times SU(3)_L \times SU(3)_R$ of E_6. The fundamental **27** representation of E_6 is then given by

$$\mathbf{27} = (3,3,1) + (3^*,1,3^*) + (1,3^*,3). \tag{3}$$

Under the decompositions $SU(3)_L \to SU(2)_L \times U(1)_{Y_L}$ and $SU(3)_R \to U(1)_{T_{3R}} \times U(1)_{Y_R}$, the individual left-handed fermionic components are defined as follows.[8]

$$\begin{aligned}
(u,d) &\sim (3;2,\frac{1}{6};0,0), && (4)\\
(\nu_e,e) &\sim (1;2,-\frac{1}{6};0,-\frac{1}{3}), && (5)\\
u^c &\sim (3^*;0,0;-\frac{1}{2},-\frac{1}{6}), && (6)\\
d^c &\sim (3^*;0,0;\frac{1}{2},-\frac{1}{6}), && (7)\\
e^c &\sim (1;0,\frac{1}{3};\frac{1}{2},\frac{1}{6}), && (8)\\
N &\sim (1;0,\frac{1}{3};-\frac{1}{2},\frac{1}{6}), && (9)\\
h &\sim (3;0,-\frac{1}{3};0,0), && (10)\\
h^c &\sim (3^*;0,0;0,\frac{1}{3}), && (11)\\
(\nu_E,E) &\sim (1;2,-\frac{1}{6};-\frac{1}{2},\frac{1}{6}), && (12)\\
(E^c,N_E^c) &\sim (1;2,-\frac{1}{6};\frac{1}{2},\frac{1}{6}), && (13)\\
S &\sim (1;0,\frac{1}{3};0,-\frac{1}{3}). && (14)
\end{aligned}$$

Note that the electric charge is given here by

$$Q = T_{3L} + Y_L + T_{3R} + Y_R, \tag{15}$$

and there are three families of these fermions and their bosonic superpartners.

Consider now the $SO(10)$ decomposition of the **27** representation:

$$\mathbf{27} = \mathbf{16} + \mathbf{10} + \mathbf{1}. \tag{16}$$

Two options are available. The conventional one (Option A) is

$$\begin{aligned} \mathbf{16} &= (u,d) + u^c + e^c + d^c + (\nu_e, e) + N, && (17) \\ \mathbf{10} &= h + (E^c, N_E^c) + h^c + (\nu_E, E), && (18) \\ \mathbf{1} &= S. && (19) \end{aligned}$$

The alternative one (Option B) is[9]

$$\begin{aligned} \mathbf{16} &= (u,d) + u^c + e^c + h^c + (\nu_E, E) + S, && (20) \\ \mathbf{10} &= h + (E^c, N_E^c) + d^c + (\nu_e, e), && (21) \\ \mathbf{1} &= N. && (22) \end{aligned}$$

The latter is obtained from the former by the exchange[9]

$$d^c \leftrightarrow h^c, \quad (\nu_e, e) \leftrightarrow (\nu_E, E), \quad N \leftrightarrow S, \tag{23}$$

so that $SU(3)_R$ is broken along a different direction, namely that given by

$$T'_{3R} = \frac{1}{2}T_{3R} + \frac{3}{2}Y_R, \quad Y'_R = \frac{1}{2}T_{3R} - \frac{1}{2}Y_R. \tag{24}$$

As far as the standard $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge symmetry is concerned, the two options are identical because

$$T'_{3R} + Y'_R = T_{3R} + Y_R = Q - T_{3L} - Y_L. \tag{25}$$

In the $U(1)_\eta$ extension, it can also be shown that there is no difference because Q_η is proportional to $T_{3L} + 5Y_L - Q$.[8] Furthermore, the same Yukawa terms are allowed by either option,

independent of any additional $U(1)$. This is easily seen by expressing the **27** representation in terms of its ($SO(10)$, $SU(5)$) components:

$$\mathbf{27} = (\mathbf{16}, \mathbf{10}) + (\mathbf{16}, \mathbf{5}^*) + (\mathbf{16}, \mathbf{1}) + (\mathbf{10}, \mathbf{5}) + (\mathbf{10}, \mathbf{5}^*) + (\mathbf{1}, \mathbf{1}). \tag{26}$$

The allowed terms must then be of the form $(\mathbf{16},\mathbf{10})(\mathbf{16},\mathbf{10})(\mathbf{10},\mathbf{5})$, $(\mathbf{16},\mathbf{10})(\mathbf{16},\mathbf{5}^*)(\mathbf{10},\mathbf{5}^*)$, $(\mathbf{10},\mathbf{5})(\mathbf{16},\mathbf{5}^*)(\mathbf{16},\mathbf{1})$, and $(\mathbf{10},\mathbf{5})(\mathbf{10},\mathbf{5}^*)(\mathbf{1},\mathbf{1})$, which remain the same if $(\mathbf{16},\mathbf{5}^*)$ and $(\mathbf{10},\mathbf{5}^*)$ are exchanged together with $(\mathbf{16},\mathbf{1})$ and $(\mathbf{1},\mathbf{1})$, in accordance with Eq. (23).

Two $U(1)$ factors are conventionally defined in Option A by the symmetry breaking chain

$$E_6 \to SO(10) \times U(1)_\psi, \quad SO(10) \to SU(5) \times U(1)_\chi. \tag{27}$$

If the extended gauge model contains only one additional $U(1)$ factor, it must be a linear combination of $U(1)_\psi$ and $U(1)_\chi$. Let

$$Q(\alpha) = Q_\psi \cos\alpha + Q_\chi \sin\alpha, \tag{28}$$

then the $U(1)_\eta$ from flux breaking corresponds to $\tan\alpha = \sqrt{3/5}$. On the other hand, the $U(1)$ factor for which N is trivial is clearly that which would be called $U(1)_\chi$ in Option B. This turns out to be given by $\tan\alpha = -\sqrt{1/15}$. To obtain this factor which will be called $U(1)_N$ from now on, the flux breaking of E_6 must be augmented by the usual Higgs mechanism, presumably at near the same scale. Consider then a pair of superheavy **27** and $\mathbf{27}^*$ representations. Assume that they develop *vev*'s along the N and N^* directions respectively. Then E_6 is broken down to the $SO(10)$ of Option B. Assume also that the flux mechanism breaks $SU(3)_L$ to $SU(2)_L \times U(1)_{Y_L}$, *i.e.* along the (1,0) direction, and $SU(3)_R$ to $U(1)_{T_{3R}} \times U(1)_{Y_R}$, *i.e.* along the (3,0) direction. Then the resulting gauge symmetry is exactly $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_N$ with

$$Q_N = 6Y_L + T_{3R} - 9Y_R. \tag{29}$$

The individual particle assignments under $U(1)_N$ are then

$$(u,d), u^c, e^c \quad : \quad 1, \tag{30}$$

$$d^c, (\nu_e, e) \quad : \quad 2, \tag{31}$$

$$h, (E^c, N_E^c) \quad : \quad -2, \tag{32}$$

$$h^c, (\nu_E, E) \quad : \quad -3, \tag{33}$$

$$S \quad : \quad 5, \tag{34}$$

$$N \quad : \quad 0. \tag{35}$$

As in any other superstring-inspired E_6 model, a discrete symmetry must be imposed to eliminate rapid proton decay.[10] Here a Z_2 symmetry is assumed where all superfields are odd except one copy each of (ν_E, E), (E^c, N_E^c), and S, which are even. Consequently, the allowed cubic terms of the superpotential are $u^c(uN_E^c - dE^c)$, $d^c(uE - d\nu_E)$, $e^c(\nu_e E - e\nu_E)$, $S(EE^c - \nu_E N_E^c)$, Shh^c, and $N(\nu_e N_E^c - eE^c)$. As the scalar components of the even superfields ν_E, N_E^c, and S acquire vev's, all particles obtain masses in the usual way. In addition, since N is now a gauge singlet, it may acquire a large Majorana mass from nonrenormalizable interactions.[7] The quadratic terms $d^c h$ and $\nu_e N_E^c - eE^c$ are also gauge singlets, and allowed by the discrete Z_2 symmetry. [The latter term is of course restricted to the two odd (E^c, N_E^c) doublets.] They are soft terms which reduce the symmetries of the Lagrangian and may thus be assumed to be naturally small.[11] Their origin is presumably also from nonrenormalizable interactions. Note that both baryon number and lepton number remain conserved.

Consider now the 5×5 mass matrix spanning ν_e, N, ν_E, N_E^c, and S. It is exactly given by combining Eq. (1) with Eq. (2) and adding a $\nu_e N_E^c$ term:

$$\mathcal{M}_5 = \begin{pmatrix} 0 & m_D & 0 & m_3 & 0 \\ m_D & m_N & 0 & 0 & 0 \\ 0 & 0 & 0 & m_E & m_1 \\ m_3 & 0 & m_E & 0 & m_2 \\ 0 & 0 & m_1 & m_2 & 0 \end{pmatrix}, \tag{36}$$

where m_D and m_1 come from $\langle \tilde{N}_E^c \rangle$, m_2 from $\langle \tilde{\nu}_E \rangle$, and m_E from $\langle \tilde{S} \rangle$. The S fermion corresonding to the last *vev* is even under Z_2 and it becomes massive because $U(1)_N$ is broken by $\langle \tilde{S} \rangle$ through which a mass term is generated, linking it with the corresponding gauge fermion. The two odd S's remain light and are naturally suited to be light singlet neutrinos. For illustration, let $m_1 = m_2 = m_e = 0.5$ MeV and $m_E = 2 \times 10^5$ GeV, then $m_S \simeq 2m_1 m_2/m_E = 2.5 \times 10^{-3}$ eV. Assuming that m_{ν_e} is much smaller, then $\nu_e - \nu_S$ oscillations occur with $\Delta m^2 \simeq 6 \times 10^{-6}$ eV^2 which is in the right range to account for the solar data. The mixing angle between ν_e and ν_S is given by $m_3/2m_2$ and should be about 0.04 for $\sin^2 2\theta \simeq 6 \times 10^{-3}$.

In conclusion, a supersymmetric extended gauge model based on $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_N$ has been proposed. This gauge symmetry is derivable from an E_6 superstring model through a combination of flux breaking and the usual Higgs mechanism with a pair of superheavy $\mathbf{27}$ and $\mathbf{27}^*$ representations. Its particle content consists of three supermultiplets belonging to the fundamental $\mathbf{27}$ representation of E_6 as listed in Eqs. (4) to (14) and Q_N is given by Eq. (29). The three N singlets are trivial under $U(1)_N$ and naturally acquire large Majorana masses from gravitationally induced nonrenormalizable interactions. One of the S singlets has a *vev* which breaks $U(1)_N$ at an unspecified scale and renders all remaining particles heavy except for those of the supersymmetric standard model and the other two S singlets. At and below the electroweak energy scale, this model differs from the minimal supersymmetric standard model (MSSM) in the following important ways. (1) The three known doublet neutrinos ν_e, ν_μ, and ν_τ have small Majorana masses instead of being massless as in the MSSM. (2) Two light singlet neutrinos exist and they may have small mixings with ν_e, ν_μ, and ν_τ. This allows for a comprehensive understanding of neutrino oscillations as well as hot dark matter in the face of all available data. (3) The scalar partners of one set of the (ν_E, E) and (E^c, N_E^c) superfields are identified with the two usual Higgs doublets Φ_1 and Φ_2

of the MSSM. However, the Higgs potentials are different because the superpotential here has the cubic term $(\nu_E N_E^c - EE^c)S$ whereas the MSSM has the quadratic term $\phi_1^0\phi_2^0 - \phi_1^-\phi_2^+$. Hence the corresponding higgsino mass is bounded here by $\langle\tilde{S}\rangle$ whereas in the MSSM, there is no understanding as to why this mass should be much smaller than the unification scale of 10^{16} GeV or the Planck scale of 10^{19} GeV. The Higgs potential of this model has only two doublets at the electroweak energy scale, but because of the above-mentioned cubic term in the superpotential, it differs from that of the MSSM by one extra coupling. Previous such examples have been given for other gauge extensions.[12]

ACKNOWLEDGEMENT

I thank A. Mendez for hospitality and a discussion at the Univ. Autonoma de Barcelona which led to this investigation. This work was supported in part by the U. S. Department of Energy under Grant No. DE-FG03-94ER40837.

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